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AUTHOR(S):

Gubler, Philipp; Kunihiro, Teiji; Lee, Su Houn

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A novel probe of chiral restoration in nuclear medium



Philipp Gubler^a, Teiji Kunihiro^b, Su Houng Lee^{a,*}

^a Department of Physics and Institute of Physics and Applied Physics, Yonsei University, Seoul 03722, Republic of Korea

^b Department of Physics, Faculty of Science, Kyoto University, Kyoto 606-8502, Japan

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ABSTRACT

We propose measuring the mass shift and width broadening of the $f_1(1285)$ meson together with those of the ω from a nuclear target as a means to experimentally probe the partial restoration of chiral symmetry inside the nuclear matter. The relation between the order parameter of chiral symmetry and the difference in the correlation functions of the $f_1(1285)$ current and the ω current is discussed in the limit where the disconnected diagrams are neglected. A QCD sum rule analysis of the $f_1(1285)$ meson mass leads to about 100 MeV attraction in nuclear matter, which can be probed in future experiments.

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1. Introduction

A long standing and prominent problem in nuclear and hadron physics is understanding the origin of the hadron masses, that comprise the dominant part of the mass of our visible universe [1–4]. While chiral symmetry breaking is expected to be responsible for generating hadron masses of the order of GeV starting from the bare quark masses that are smaller than 10 MeV, it is not clear how the effect is manifested inside a hadron, where confinement makes the quark an unobservable gauge dependent object.

There is a well defined order parameter of chiral symmetry breaking: the chiral condensate $\langle \bar{q}q \rangle$. However it is a challenge to relate it to physically accessible quantities in a model independent way [5]. The differences between current–current correlation functions of chiral partners are considered to be another promising set of order parameters of chiral symmetry. In fact, it was shown that the density of zero eigenvalues of the Dirac equation, responsible for generating the chiral order parameter [6], is also responsible for breaking the degeneracies of the corresponding correlation functions [7] apart from the $U_A(1)$ breaking effect coming from contributions of topologically non-trivial configurations [8].

While the whole correlation function should be measured to verify the restoration of the chiral symmetry breaking, the ground state poles that couple to the correlation functions are the most distinctive feature of the correlation functions, and hence measuring the mass differences between chiral partners is considered to be the most attractive alternative. Moreover, chiral symmetry

is expected to be restored in the spectrum of excited states [9]. The first works on the temperature dependence of the sigma-pion mass [1,10] showed that the masses of the chiral partners become degenerate near the chiral phase transition point. The vector–axialvector masses [11] were also found to decrease at finite temperature. One could think of measuring such mass shift of hadrons in heavy ion collisions. However, whatever signal there is near the phase transition, it will be lost during the hadronic evolution of the system. In fact, signals appearing as peaks in the low mass dilepton spectrum were found to be dominated by a large broadening of the ρ meson [12].

Measuring the mass shift from the nuclear target was suggested as an alternative because the chiral order parameter can be shown to be quenched by more than 30% in nuclear medium. Moreover, the nuclear target would provide a stable environment where the density profile is fixed so that the effects of time evolution can be neglected [4]. Unfortunately, so far, most attempts were problematic because the width broadening of the already wide mesons made any realistic measurement not meaningful. However, the CBELSA/TAPS collaboration have been focusing on the ω and η' meson and successfully measured the mass and with broadening of these particles [13,14]. The main reason behind their success is in their focus on mesons with small vacuum width so that a non-trivial proportional increase in the meson's width will still leave the meson narrow enough for a meaningful measurement to identify and describe the meson in the nuclear medium.

Recently, the CLAS collaboration has successfully identified the $f_1(1285)$ in photoproduction from a proton target with a small width of 18 ± 1.4 MeV [21]. As we will discuss in the next section, the $f_1(1285)$ can be considered to be the chiral partner of

* Corresponding author.

E-mail address: suhoung@yonsei.ac.kr (S.H. Lee).

the ω if disconnected diagrams are neglected. Hence, when the experiment is applied to a nuclear target and when combined with observations for the ω meson, one can finally hope to observe and experimentally verify the partial chiral symmetry restoration expected to occur in nuclear medium. Such an observation will provide critical information on the origin of hadron masses.

2. Ideal mixing in the vector and axial vector channels

When we restrict ourselves to two flavors, the chiral partners of $SU(2)_L \times SU(2)_R$ in the vector–axial vector channel are the ρ and a_1 mesons. Both have large width even in vacuum and are hence not good candidates to be measured in nuclear medium. Extending the flavor number to three, one finds that the octet meson is mixed with the singlet, forming a nonet. For the vector channel the extra mesons with isospin zero mixes with the octet almost ideally. This means that one can identify the quark content as $\omega = \frac{1}{\sqrt{2}}(\bar{u}u + \bar{d}d)$ and $\phi = \bar{s}s$. At the same time, the ω mass is almost degenerate with that of the ρ , which comes about naturally in the large N_c limit where we can neglect the correlation function between u -quark current and the d quark current.

In fact, ideal mixing in the three flavor case and the suppression of disconnected diagrams are related. To see this we recapitulate here briefly why we have ideal mixing in the vector channel, where the suppression of disconnected diagrams is largest. To visualize, it is convenient to just consider the two point function composed of the singlet $\omega_1 = \frac{1}{\sqrt{3}}(\bar{u}u + \bar{d}d + \bar{s}s)$ and the octet $\omega_8 = \frac{1}{\sqrt{6}}(\bar{u}u + \bar{d}d - 2\bar{s}s)$ currents. The correlation functions can be represented as a two by two matrix composed of the elements $\Pi_{ij} = \langle \omega_i, \omega_j \rangle$. If we neglect the disconnected diagrams and work in the quenched approximation, this matrix can be written as follows

$$\begin{pmatrix} \Pi_0 + \frac{1}{3}\Delta\Pi_s & -\frac{\sqrt{2}}{3}\Delta\Pi_s \\ -\frac{\sqrt{2}}{3}\Delta\Pi_s & \Pi_0 + \frac{2}{3}\Delta\Pi_s \end{pmatrix}, \quad (1)$$

where

$$\Pi_0 = \langle \omega_1, \omega_1 \rangle_{m_s=0}, \quad (2)$$

$$\Delta\Pi_s = \langle \bar{s}s, \bar{s}s \rangle - \langle \bar{s}s, \bar{s}s \rangle_{m_s=0}. \quad (3)$$

That is, Π_0 represents the $SU(3)$ symmetric correlation function, while the $\Delta\Pi_s$ encodes the symmetry breaking part and the mixing between the singlet and octet. Diagonalizing, we find the ideally mixed correlation. If we include the disconnected diagram, we will not obtain ideal mixing. Therefore, whether we are in two or three flavors, when disconnected diagrams are neglected, the ω becomes degenerate with the ρ .

In the axial vector channel, the mixing angle between the two isospin zero mesons $f_1(1285)$, $f_1(1420)$ is not determined as well as in the case of the vector mesons. Some hadronic models in fact obtain a rather large strangeness component of the $f_1(1285)$ [15]. Most estimates however find that the quark content of the $f_1(1285)$ is dominated by u and d quarks with only a small s component [16–18]. Our own analysis based on QCD sum rules, to be described in the next section, points to the same conclusion. Moreover, one should note that the mass of $f_1(1285)$ is almost degenerate with that of the a_1 , as the ω is with the ρ (Table 1). All this suggests that the $f_1(1285)$ and $f_1(1420)$ are almost ideally mixed, demonstrating that the disconnected diagrams are also suppressed in this channel. Within this limit, one can argue that the ω and the $f_1(1285)$ are chiral partners and will become degenerate when the chiral symmetry gets restored, which was also anticipated in Ref. [19].

Table 1

Width and mass of chiral partners. Units are in MeV.

$J^{PC} = 1^{--}$	Mass	Width	$J^{PC} = 1^{++}$	Mass	Width
ρ	770	150	a_1	1260	250–600
ω	782	8.49	f_1	1285	24.2
ϕ	1020	4.266	f_1	1420	54.9

Let us discuss this point in some more detail. In other words: why can ω and $f_1(1285)$ be considered as chiral partners when disconnected diagrams are neglected, while they do not seem to be directly related by chiral $SU(2)$ symmetry? This is related to why the huge ω or $f_1(1285)$ masses, relative to the current quark masses, are almost degenerate to that of the ρ and a_1 respectively. To understand this, one notes that using the Casher Banks formula, one finds that spontaneous chiral symmetry breaking occurs when the density of zero modes of the Dirac equation within the QCD Euclidean functional integral becomes non-zero. Once this density of zero modes becomes non-zero, one can show that all the order parameters of chiral symmetry breaking become non-zero. One can furthermore show that when one neglects the disconnected diagrams, the difference between the correlation functions of ω and $f_1(1285)$ is proportional to the density of zero modes [7,8,20], and are in fact identical to that of ρ and a_1 . Now, even if one includes the disconnected diagrams, one notes that their contributions are small in the vacuum as the leading contribution requires at least three gluon exchange. Moreover, the effect from the $U_A(1)$ anomaly is small in the vacuum. This is so because the phenomenologically observable effect of the $U_A(1)$ anomaly comes in through topologically non-trivial configurations that connects zero modes of different chirality. Such configurations have minimal effects in the vector or axial channel. One can visualize such effects as an instanton contribution that has left handed and right handed quarks with N_F flavors interpolating between correlation functions of the two currents; if the currents are vector or axial vector currents, a single instanton can not interpolate the correlation function as left right symmetry is conserved in such currents.

In nuclear matter, the non-strange chiral condensate $\langle \bar{u}u + \bar{d}d \rangle$ is expected to be partially quenched. From a theoretical point of view, this originates from the identical reduction of the zero-modes of the u and d quarks. On the other hand, if one still neglects disconnected diagrams, ideal mixing of ω and $f_1(1285)$ will not be modified even in nuclear matter as any additional correlations between strange and non-strange components vanish in this limit. This can be understood from a large N_c argument, as in the large N_c limit the nucleon has no sea quarks and therefore no strangeness content. Next, considering only u and d quarks and assuming that the correlation between disconnected quark lines within a nucleon is suppressed compared to those that are connected, one concludes that the changes of the zero modes will affect the difference between the ω and the $f_1(1285)$ in the same way as it does for the ρ and a_1 , when disconnected diagrams are neglected. Any residual $U_A(1)$ anomaly breaking effect in the presence of nucleons can in principle be estimated separately by measuring the η' in nuclear matter, which has been made possible by the experimental efforts reported in Ref. [14]. Therefore when the measurements of all the small width mesons, η' , ω and $f_1(1285)$ are combined, we will finally have a better understanding of how, if at all, chiral symmetry and $U_A(1)$ are partially restored in nuclear medium, and whether they are at all responsible for generating hadron masses.

3. QCD sum rule analysis

We first give a quick overview of the QCD sum rule analysis method adapted in this work.

The starting point is the correlation function of the axial vector current in the nuclear medium,

$$\Pi_{\mu\nu}(\omega, \mathbf{q}) = i \int d^4x e^{iqx} \langle T J_\mu(x) J_\nu(0) \rangle_{n.m.}, \quad (4)$$

where the subscript *n.m.* stands for the nuclear matter expectation value and $q_\mu = (\omega, \mathbf{q})$. The current is taken to be either $J_\mu^q = \eta_{\mu\nu} \frac{1}{\sqrt{2}} (\bar{u} \gamma_\nu \gamma_5 u + \bar{d} \gamma_\nu \gamma_5 d)$ or $J_\mu^s = \eta_{\mu\nu} (\bar{s} \gamma_\nu \gamma_5 s)$, where $\eta_{\mu\nu} = q_\mu q_\nu / q^2 - g_{\mu\nu}$. We will closely follow the finite temperature formalism given in Ref. [11] and apply it to finite density [3]. We look at the trace part of the polarization function $\Pi(Q^2) = -\Pi_\mu^\mu(Q^2)/3$ at $\mathbf{q} \rightarrow 0$. The OPE for J_μ^q is given as

$$\begin{aligned} \Pi(Q^2) = & \frac{1}{4\pi^2} \left(1 + \frac{\alpha_s}{\pi} \right) Q^2 \ln(Q^2) + \frac{3}{2\pi^2} m_q^2 \ln(Q^2) \\ & + 2 \frac{m_q}{Q^2} \langle \bar{q}q \rangle - \frac{1}{12Q^2} \left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle + \frac{2\pi\alpha_s}{Q^4} \langle (\bar{q} \gamma_\mu \lambda^a q)^2 \rangle \\ & + \frac{4\pi\alpha_s}{9Q^4} \langle (\bar{q} \gamma_\mu \lambda^a q) \left(\sum_q^{u,d,s} \bar{q} \gamma_\mu \lambda^a q \right) \rangle \\ & + \frac{8iq^\mu q^\nu}{3Q^4} \langle (\bar{q} \gamma_\mu D_\nu q)_{ST} \rangle \\ & - \frac{32iq^\mu q^\nu q^\lambda q^\sigma}{3Q^8} \langle (\bar{q} \gamma_\mu D_\nu D_\lambda D_\sigma q)_{ST} \rangle. \end{aligned} \quad (5)$$

Here, we take the quark operators $\langle \bar{q}..q \rangle$ to be the average of *u, d* quark contributions, except for the one with the summation sign. To get the respective expression for the J_μ^s correlator, one simply replaces $\langle \bar{q}..q \rangle$ by $\langle \bar{s}..s \rangle$, again with the exception of the operator behind the summation sign. While for *u* and *d* quarks, the m_q^2 term in Eq. (5) is negligible and can be safely neglected, we keep it for the strange quark case. We have also neglected the twist-4 operators in Eq. (5), whose contributions are expected to be small [22].

We now follow the standard procedure and use the dispersion relation to relate the OPE to the spectral density, generally consisting of poles and continuum, which we represent by a delta function for the lowest pole and a step function starting from a continuum threshold s_0 , respectively; in the present simplest treatment, the width of the pole is neglected. Performing the Borel transformation and taking the ratio with its derivative, one obtains a relation for the mass in terms of the Borel mass:

$$\begin{aligned} \frac{m_{f_1}^2}{M^2} = & \left[2 \left(1 + \frac{\alpha_s}{\pi} \right) E_2(s_0/M^2) - \frac{a}{M^2} E_1(s_0/M^2) \right. \\ & \left. + \frac{2(e-f)}{M^6} \right] \\ & \times \left[\left(1 + \frac{\alpha_s}{\pi} \right) E_1(s_0/M^2) - \frac{a}{M^2} E_0(s_0/M^2) \right. \\ & \left. - \frac{-b+c+d}{M^4} - \frac{2(e-f)}{M^6} \right]^{-1}. \end{aligned} \quad (6)$$

In the present sum rule, there is no contribution from the nucleon scattering term [22]. This is so because the scattering term contributes as $s\delta(s)$ in the imaginary part of $\Pi(Q^2)$. If we had studied the sum rule for $\Pi(Q^2)/Q^2$, the scattering term would have appeared not only for the f_1 , but also in the ω , ρ and a_1 sum rules with different coefficients. To linear order in density, additional scattering terms coming from excited nucleon intermediate states

could however be added to the imaginary part. Such terms will not be of the delta function type but appear near the excitation energies for the intermediate states and will depend on the quantum numbers of the current. We leave such a detailed modeling as a future work. Here we follow the simple pole ansatz to estimate the maximum possible mass shift for the f_1 meson and compare it to the vacuum width which is an important criterion to assess the observability in an actual experiment. The parameters *a–f* read

$$a = 6m_q^2, \quad (7)$$

$$b = 4\pi^2 m_q \langle \bar{q}q \rangle_\rho = 4\pi^2 (m_q \langle \bar{q}q \rangle_0 + \sigma_{\pi N} \rho), \quad (8)$$

$$c = \frac{\pi^2}{3} \left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle_\rho \quad (9)$$

$$= \frac{\pi^2}{3} \left(\left\langle \frac{\alpha_s}{\pi} G^2 \right\rangle_0 - \frac{8}{9} (M_N - \sigma_{\pi N} - \sigma_{sN}) \rho \right), \quad (10)$$

$$d = 4\pi^2 M_N A_2^q \rho, \quad (11)$$

$$e = \frac{704\pi^3 \alpha_s}{4 \times 81} \langle \bar{q}q \rangle_\rho^2, \quad (12)$$

$$f = \frac{10\pi^2}{3} M_N^3 A_4^q \rho, \quad (13)$$

and the functions $E_i(s_0/M^2)$ are defined as

$$E_0(s_0/M^2) = 1 - e^{-s_0/M^2}, \quad (14)$$

$$E_1(s_0/M^2) = 1 - (1 + \frac{s_0}{M^2}) e^{-s_0/M^2}, \quad (15)$$

$$E_2(s_0/M^2) = 1 - (1 + \frac{s_0}{M^2} + \frac{s_0^2}{2M^4}) e^{-s_0/M^2}. \quad (16)$$

The threshold parameter is determined by requiring the mass in Eq. (6) to be most stable within the Borel window. The minimum Borel mass is determined by requiring the contribution coming from the highest OPE to be less than 10% of the whole OPE appearing in the denominator of Eq. (6). The maximum Borel mass is determined by the pole dominance criterion, demanding that the pole contribution to the sum rule is larger than that of the continuum.

3.1. Vacuum analysis of $f_1(1285)$ and $f_1(1420)$

At first, let us study the masses extracted from the J_μ^q and J_μ^s correlators to check whether they can approximately reproduce the energy levels of the physical $f_1(1285)$ and $f_1(1420)$ states. If J_μ^q (J_μ^s) couples dominantly to $f_1(1285)$ ($f_1(1420)$), this would be a strong indication that these two states are close to ideally mixed.

The Borel curves for the vacuum masses of the J_μ^q and J_μ^s correlators are shown in Fig. 1. The parameter values employed to draw these curves are given in the upper part of Table 2. The obtained curves are stable and exhibit a wide Borel window, which indicates that the sum rule analysis works well for both cases. The flattest Borel curves give threshold parameters of $\sqrt{s_0} = 1.61$ GeV for J_μ^q and $\sqrt{s_0} = 1.925$ GeV for J_μ^s . Comparing the calculated masses with the experimental values, it is seen that the result from the J_μ^q correlator agrees almost perfectly with the mass of the $f_1(1285)$. The curve extracted from J_μ^s similarly lies close to the $f_1(1420)$ mass, however turns out to be about 70 MeV too high. Given the slight discrepancy existing in the $f_1(1420)$ sum rule, further theoretical studies are required to shed more light on the possible importance of disconnected diagrams in the f_1 channel and substantiate the chiral partner scenario. Furthermore, there could still be a significant four-quark component in the wave functions for both $f_1(1285)$ and $f_1(1420)$, as has been shown for instance in

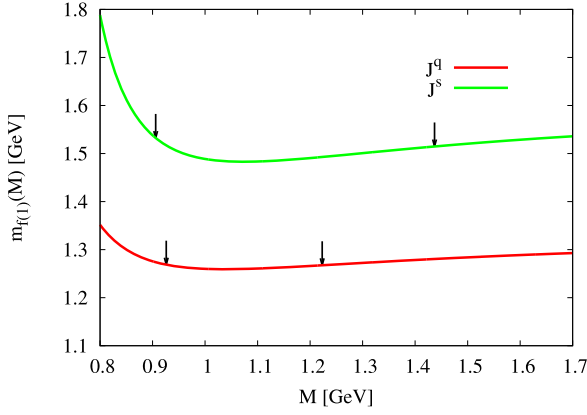


Fig. 1. The Borel curve for the vacuum mass formula of Eq. (3) for the J_μ^q and J_μ^s correlators. The arrows indicate the positions of the respective minimum and maximum Borel masses.

Table 2

Parameter values used in the present calculation. All values were converted to a renormalization scale of 1 GeV. The parameters A_2^q and A_4^q are obtained by numerically integrating the NLO parton distributions provided in [32].

$\langle \bar{q}q \rangle_0$	$(-0.248 \text{ GeV})^3$ [26]
$\langle \bar{s}s \rangle_0$	$0.8 \times \langle \bar{q}q \rangle_0$ [27]
m_q	4.7 MeV [28]
m_s	95 MeV [28]
$\langle \frac{\alpha_s}{\pi} G^2 \rangle_0$	0.012 GeV^4 [29]
M_N	939 MeV
$\sigma_{\pi N}$	$45 \pm 15 \text{ MeV}$ [24,25,30]
σ_{sN}	35 MeV [31]
A_2^q	0.62 [32]
A_4^q	0.066 [32]

the coupled channel type analysis of Ref. [23], which contributes to both the connected and disconnected diagrams. Nevertheless, the sum rule results indeed show that in the vacuum the J_μ^q and J_μ^s currents couple strongly to the $f_1(1285)$ and $f_1(1420)$, respectively, and that therefore these two states are mixed almost ideally.

3.2. Finite density analysis of $f_1(1285)$

Let us next turn to the main topic of this paper: the modification of the $f_1(1285)$ at finite density. This QCD sum rule analysis should provide a guideline for the expected mass shift of the $f_1(1285)$ in nuclear matter.

The parameters used to quantify the density dependence of the various condensates are given in the lower part of Table 2. The corresponding result is shown in Fig. 2, where the Borel curves for both the vacuum and normal nuclear matter density ρ_0 are plotted. The flattest Borel curve in the latter case was obtained for a threshold parameter of $\sqrt{s_0} = 1.49 \text{ GeV}$.

As it is seen in Table 2, the analysis was performed with a central value of 45 MeV for the πN sigma term $\sigma_{\pi N}$. However, there are lattice results that show that $\sigma_{\pi N}$ might be smaller [24] while a recent phenomenological fit to experimental πN scattering data suggests that it is bigger [25]. We have therefore performed the analysis for $\sigma_{\pi N} = 30 \text{ MeV}$ and 60 MeV to check the sensitivity of our results on the πN sigma term value, which is the largest source of error for the f_1 mass shift δm_{f_1} and leads to the bands shown in Fig. 3, which depicts δm_{f_1} as a function of density. Taking this uncertainty into account, we expect a mass shift of about $96 \pm 38 \text{ MeV}$ in nuclear medium.

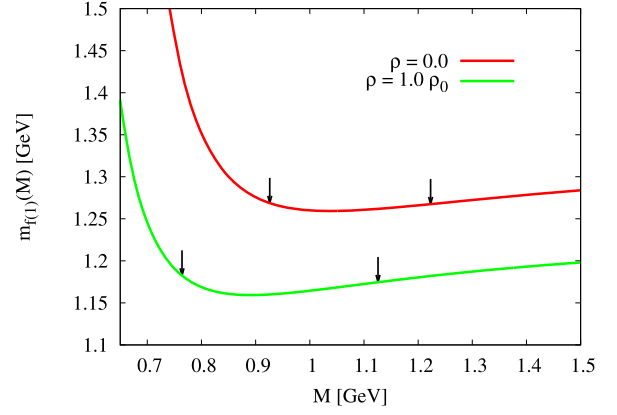


Fig. 2. The Borel curve for the mass in Eq. (3) in the vacuum and at nuclear matter density. The arrows show the position of the minimum and maximum Borel mass.

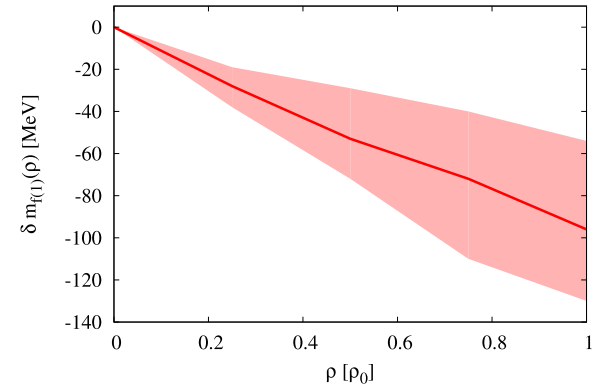


Fig. 3. The expected density dependence of the $f_1(1285)$ meson mass shift. The solid line is obtained with a value of 45 MeV for the πN sigma term. The lower and upper bounds were respectively obtained with $\sigma_{\pi N} = 60 \text{ MeV}$ and 30 MeV .

4. Discussion and conclusion

The mass shift in the sum rule is obtained by assuming a delta function pole for the $f_1(1285)$ in nuclear medium. It was however noted in an earlier work that for the vector meson sum rule, the changes of the OPE in the medium can also be satisfied with a smaller change in the mass and a simultaneous increase in the width [33]. It is likely that a similar effect also applies to the $f_1(1285)$ discussed in this paper. Therefore, our result should be considered as a maximum mass shift value expected at nuclear matter. The experimental result for the ω suggests a small mass shift of -29 MeV and a larger increase in the width of 70 MeV [13]. Once the medium modification for the $f_1(1285)$ is experimentally observed, one can construct QCD sum rules for the ω and $f_1(1285)$ separately and analyze how the changes in the corresponding masses and widths are related to the changes in the condensates. From such analysis, one can then also estimate the effects of the factorizable part of the four quark condensate to the properties of the ω and $f_1(1285)$ meson. We leave such a detailed QCD sum rule analysis for both the ω and $f_1(1285)$ meson as future work.

The CLAS collaboration was able to clearly identify a sharp peak for the $f_1(1285)$ on a proton target [21]. Performing the experiment on a nuclear target will involve several difficulties. First of all, the present mass shift is obtained with the $f_1(1285)$ meson at rest with respect to the nuclear medium. Experimentally selecting out low momentum $f_1(1285)$ will strongly suppress the signal. Moreover, reconstructing the $f_1(1285)$ from the hadronic final states

will entail smearing and/or lost signal due to the rescattering of the final state with the medium.

However, as we have emphasized in this work, within the limit where disconnected diagrams can be neglected, such a measurement would be the first direct observation of a chiral symmetry restoration effect on hadron properties. Based on theoretical estimates on how much the chiral order parameter would change at finite density, chiral symmetry is expected to be partially restored in nuclear matter. Experimentally observing the $f_1(1285)$ in nuclear matter would therefore serve as a test for these theoretical expectations and hence could shed light into the mechanism of how the mass of hadrons are generated. Considering the reward, the difficulties are worth overcoming.

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